Photoexcitation of the Nucleon in the Constituent Quark Model

I. P. Matamba([1]), M. Braun([2]) and S. A. Sofianos([2])

([1]) *Physics Department, University of Venda, Private Bag X5050, Thohoyandou 0950, South Africa*

([2]) *Physics Department, University of South Africa, P.O.Box 392, Pretoria 0003, South Africa*

(ricevuto ?; approvato ?)

Summary. — Photoexcitation of the nucleon is studied in the framework of the constituent quark model with nonrelativistic quark-quark interactions. Utilizing the dominance of the pairwise structure of the relevant three-quark potential, the bound and excited state wave functions have been constructed using the Faddeev-type Integrodifferential Equation Approach. The method is based on the Hyperspherical Harmonics expansion and takes two-body correlations exactly into account and therefore reliable wave functions, required in photo processes, can be obtained. The integrated photoabsorption cross sections for the $E1$ and $M1$ transitions are calculated and compared with experimental values as well as with those of other theoretical calculations. Our results, in overall, are in good to very good agreement with the experimental values.

PACS 12.39 – Quark models.
PACS 12.39.Pn – Potential models.
PACS 14.20.Gk – Baryon resonances.

1. – Introduction

In the nonrelativistic constituent quark model for mesons and baryons the quarks are bound by confining potentials. Despite their limitations and concerns about their validity, these potential models describe the properties of the various mesons and baryons surprisingly well (see, for example, Silvestre-Brac [1] and the two excellent review articles by Lucha et al. [2] and Richard [3] on the matter).

Once the potential model approach is adopted, the three-quark problem can be solved using various approaches. Among them, the Hyperspherical Harmonics (HH) method is quite successful in applications as it is well suited to describe radial and orbital excitations [4, 5, 6, 7]. Within the framework of the HH method, the Hypercentral Approximation (HCA) has been used in the past [8, 9, 10] to study the spectrum of the baryon. There are various reasons for adopting the HCA to study the three quark system: i) The two-body potential acting between quarks is quite soft and therefore in the HH expansion of the interaction only the first term gives a significant contribution to the

binding energy. This of course means that the two-body correlations are not as strong as compared to the nuclear correlations; ii) it is quite simple and thus one avoids the complicated three-body calculations via, for example, the Faddeev equations [1, 11], and iii) the results obtained from it are accurate and the spectra are well reproduced.

Another method, in the framework of the HH method, is the Integrodifferential Equation Approach (IDEA) [12, 13, 14, 15] which includes higher terms of the HH expansion in an average manner. The IDEA method takes two-body correlations into account exactly, reproduces the spectrum of the nucleon quite well, and provides wave functions reliably [14, 15] which is crucial in studying photoexcitation processes. These processes are manifested as resonances and can be excited through electromagnetic transitions giving rise to large enhancements in the total absorption cross section [16].

The photoexcitation of the nucleon resonances has been studied in the past by various groups [16, 17, 18]. The results obtained by them are rather unsatisfactory when compared to the experimental data. The inclusion of retardation effects and relativistic corrections does not improve the situation much [16, 17]. In this work we consider the absorption of a single photon by a nucleon which then undergoes a transition from the ground state to an excited one. The photoabsorption cross section is calculated using various quark-quark potentials and by using the HCA and IDEA methods.

In Sec. 2 we describe our formalism. In Sec. 3 we give details on how the $E1$ and $M1$ transition amplitudes are calculated while in Sec. 4 we present our results and discussions.

2. – Formalism

The photoexcitation process is described by the transition amplitude

$$M_{fi} =< \Psi_f|H|\Psi_i > , \tag{1}$$

where Ψ_i is the initial ground state wave function of the nucleon, Ψ_f is the wave function of the final excited state, and H the perturbative electromagnetic Hamiltonian. In what follows we shall discuss these ingredients in some detail.

2˙1. *The wave functions.* – The fully antisymmetric total wave function for a three-quark system can be expressed as a product of configuration space, flavor, spin, and color functions. Since baryons are color singlets, the color wave function is totally antisymmetric (A) and thus the remaining product must be fully symmetric (S),

$$\Psi^A_{\rm total} = \underbrace{\psi_{\rm space} \times \Phi_{\rm flavor} \times \chi_{\rm spin}}_{S} \times \underbrace{C_{\rm color}}_{A} . \tag{2}$$

The structure of the symmetric component of Eq. (2) depends on the transition considered and can be constructed using the various symmetries involved.

For the construction of the symmetric part of the total wave function the fully symmetric, mixed symmetric, and mixed antisymmetric configuration space wave functions are required. These can be obtained using the IDEA [12, 13] method. In this method the fully symmetric ground state configuration space wave function is constructed from the Faddeev-type components $P(z, r)$ [15]

$$\Psi^S(\vec{\rho}, \vec{\sigma}) = \frac{1}{r^{5/2}} \left[P^S(z_{12}, r) + P^S(z_{23}, r) + P^S(z_{31}, r)\right] , \tag{3}$$

where $(\vec{\rho},\vec{\sigma})$ are the Jacobi coordinates, $r=\left[\frac{2}{3}\sum_\alpha \rho_\alpha^2\right]^{1/2}$ is the hyperradius with $\rho_\alpha = r_\alpha\ , \alpha = 12, 23, 31$, and the z_α are given by $z_\alpha = 2\rho_\alpha^2/r^2 - 1$. The required mixed symmetry states for $L=1$ are given by

$$\begin{aligned}\Psi_1^{M^S}(\vec{\rho},\vec{\sigma}) &= \frac{1}{r^{5/2}}\Bigg\{(1+z_{12})^{1/2}Y_{10}(\omega_{12})P_1^{S'}(z_{12},r)\\ &-\frac{1}{2}\Bigg[(1+z_{23})^{1/2}Y_{10}(\omega_{23})P_1^{S'}(z_{23},r)\\ &+(1+z_{31})^{1/2}Y_{10}(\omega_{31})P_1^{S'}(z_{31},r)\Bigg]\Bigg\}, \end{aligned} \tag{4}$$

$$\begin{aligned}\Psi_1^{M^A}(\vec{\rho},\vec{\sigma}) &= \frac{1}{r^{5/2}}\Bigg[(1+z_{31})^{1/2}Y_{10}(\omega_{31})P_1^{S'}(z_{31},r)\\ &-(1+z_{23})^{1/2}Y_{10}(\omega_{23})P_1^{S'}(z_{23},r)\Bigg], \end{aligned} \tag{5}$$

where the superscripts M^S and M^A denote the mixed symmetric and antisymmetric states with respect to the interchange of particles 1 and 2.

The required symmetric spin-flavor states are given by

$$\left|\xi^S\right\rangle = \frac{1}{\sqrt{2}}\left[\Phi^{M^S}\chi^{M^S}+\Phi^{M^A}\chi^{M^A}\right], \tag{6}$$

while the mixed symmetric states are

$$\left|\xi^{M^S}\right\rangle = \frac{1}{\sqrt{2}}\left[\Phi^{M^S}\chi^{M^S}-\Phi^{M^A}\chi^{M^A}\right], \tag{7}$$

$$\left|\xi^{M^A}\right\rangle = \frac{1}{\sqrt{2}}\left[\Phi^{M^S}\chi^{M^A}+\Phi^{M^A}\chi^{M^S}\right]. \tag{8}$$

The relevant flavor and spin states are given by various authors and therefore, will not be presented here (see, for example, Refs. [18, 19]). The singlet, antisymmetric color state,

$$C^A_{color} = \frac{1}{\sqrt{6}}(RBY - BRY + BYR - YBR + YRB - RYB), \tag{9}$$

where R,B and Y stand for Red, Blue, and Yellow respectively, does not enter into the calculations and therefore, in what follows will be suppressed.

The initial total wave function for the proton (P) ground state, with $L=0$, $S=1/2$, and $J=1/2$, is given by

$$|\Psi_i\rangle = \frac{1}{\sqrt{2}}\left[\Phi_{\rm P}^{M^S}\chi_{\rm P}^{M^S}+\Phi_{\rm P}^{M^A}\chi_{\rm P}^{M^A}\right]\left|\Psi_0^S\right\rangle, \tag{10}$$

where the lower index of the space wave function $|\Psi_0^S>$ refers to the angular momentum L. The final wave function for the first excited state, with $L=1$, $S=1/2$, and $J=1/2$ or $3/2$, of the proton is

$$\left|\Psi_f\right\rangle = \frac{1}{2}\left[\left(\Phi_{\rm P}^{M^S}\chi_{\rm P}^{M^S}-\Phi_{\rm P}^{M^A}\chi_{\rm P}^{M^A}\right)\left|\Psi_1^{M^S}\right\rangle+\left(\Phi_{\rm P}^{M^S}\chi_{\rm P}^{M^A}+\Phi_{\rm P}^{M^A}\chi_{\rm P}^{M^S}\right)\left|\Psi_1^{M^A}\right\rangle\right]. \tag{11}$$

For the $M1$ transition $(S = 1/2) \rightarrow (S = 3/2)$, where the proton and $\Delta^+(1232)$ both have an angular momentum $L = 0$, the total wave function for the initial state of the proton is still given by Eq. (10), while the final wave function for the delta is

$$\left|\Psi_f\right\rangle = \Phi^S_\Delta \chi^S_\Delta \left|\Psi^S_\Delta\right\rangle \,. \tag{12}$$

2˙2. *Electromagnetic transitions.* – The perturbative Hamiltonian for the electric dipole E1 transition, in the case of three quarks of equal mass m, is given by

$$H_{E1} = -\frac{1}{mc}\sum_{j=1}^{3} \lambda_j \hat{\epsilon}_\gamma \cdot \vec{p}_j \,, \tag{13}$$

where $\vec{p}_j$ is the momentum of quark j and $\hat{\epsilon}_\gamma$ denotes a polarization direction of the incident photon. For u and d quarks, the charge operator has the form

$$\lambda_j = \frac{e}{6}(1 + 3\tau^z_j) \,, \tag{14}$$

where τ^z_j is the third component of the isospin of the j-th quark with

$$\tau^z_j|u> = |u>, \qquad \tau^z_j|d> = -|d> \,. \tag{15}$$

Using commutation relations to express the momenta in terms of the three-quark Hamiltonian [20] we may rewrite (13) in a Siegert form [21, 22]

$$H_{E1} = -\frac{i}{\hbar c}(E_f - E_i)\sum_{j=1}^{3} \lambda_j \hat{\epsilon}_\gamma \cdot \vec{x}_j \,, \tag{16}$$

where $\vec{x}_j$ is the coordinate of j-th quark conjugate to the momentum $\vec{p}_j$. In Jacobi coordinates we have

$$H_{E1} = -\frac{1}{2i\hbar c}(E_f - E_i)(\hat{\epsilon}_\gamma \cdot \vec{\rho})I_p - \frac{\sqrt{3}}{3i\hbar c}(E_f - E_i)(\hat{\epsilon}_\gamma \cdot \vec{\sigma})I_q \,, \tag{17}$$

where $(E_f - E_i)$ is the difference between the final and initial binding energies and the operators Ip and I_q are given by

$$\begin{aligned} I_p &= \frac{e}{2}(\tau^z_1 - \tau^z_2), \\ I_q &= \frac{e}{2}\left(\frac{\tau^z_1 + \tau^z_2}{2} - \tau^z_3\right) \,. \end{aligned} \tag{18}$$

Thus, instead of expressing the Hamiltonian in terms of the individual particle charge and coordinates, the more appropriate Jacobi coordinates and operators I_p and I_q which act on quasi-particles [23] are used.

The magnetic dipole M_1 causes the transition $\gamma\mathrm{P} \rightarrow \Delta^+(1232)$, in which the proton, after absorbing a photon (γ), is excited to the delta (Δ^+). The corresponding perturbative Hamiltonian is

$$H_{M1} = -i\sum_{j=1}^{3}\left(\vec{\mu}^j_q \times \vec{k_\gamma}\right) \cdot \hat{\epsilon}_\gamma \,, \tag{19}$$

where $\vec{\mu}^j_q$ is the magnetic moment operator of the j-th quark. Since H_{M1} does not contain any orbital operators, in this transition the spin must change instead.

3. – The transition amplitudes

Noting that $E_\gamma = E_f - E_i$ and by letting the charge operators I_p and I_q act on the three-quark isospin states we finally obtain for the transition matrix elements

$$M_{E1} = \frac{eE_\gamma}{2\sqrt{6}i\hbar c}\left(\left\langle \Psi_1^{M^A} \left| \hat{\epsilon}_\gamma \cdot \vec{\rho} \right| \Psi_0^S \right\rangle - \left\langle \Psi_1^{M^S} \left| \hat{\epsilon}_\gamma \cdot \vec{\sigma} \right| \Psi_0^S \right\rangle \right). \tag{20}$$

The integrals in (20) were evaluated using Euler angles α, β, γ as external and ρ, σ, $x = \vec{\rho}\cdot\vec{\sigma}/\rho\sigma$ as internal coordinates. Here the z' axis is chosen to coincide with $\hat{\rho}$ and $\hat{\sigma}$ is lying in the $x' - z'$ plane. Thus only a five dimensional integration has to be done numerically since both Ψ_0^S and the two components of Ψ_1 are invariant with respect to a rotation about the z-axis and thus do *not* depend on α. After averaging over the direction of $\vec{k}$ and polarization direction one obtains the following expression for the absolute square of the transition matrix elements

$$\overline{\left|\mathcal{M}_{E1}\right|^2} = \frac{e^2 E_\gamma^2}{72\,(\hbar c)^2}\left| \left\langle \Psi_1^{M^A} \middle| \rho_z \middle| \Psi_0^S \right\rangle - \left\langle \Psi_1^{M^S} \middle| \sigma_z \middle| \Psi_0^S \right\rangle \right|^2. \tag{21}$$

Therefore, the following integrals are required

$$\begin{aligned} \left\langle \Psi_1^{M^A} \middle| \rho_z \middle| \Psi_0^S \right\rangle = 2\pi \int_0^\pi \sin\beta d\beta \int_0^{2\pi} d\gamma \int_0^\infty \rho^2 d\rho \int_0^\infty \sigma^2 d\sigma \int_{-1}^1 dx \\ \Psi_1^{M^A}(\rho,\sigma,x,\beta,\gamma)\rho\cos\beta\,\Psi_0^S(\rho,\sigma,x)\,, \end{aligned} \tag{22}$$

$$\begin{aligned} \left\langle \Psi_1^{M^S} \middle| \sigma_z \middle| \Psi_0^S \right\rangle = 2\pi \int_0^\pi \sin\beta d\beta \int_0^{2\pi} d\gamma \int_0^\infty \rho^2 d\rho \int_0^\infty \sigma^2 d\sigma \int_{-1}^1 dx \\ \Psi_1^{M^S}(\rho,\sigma,x,\beta,\gamma)\sigma(\cos\beta\cos\theta - \sin\beta\cos\gamma\sin\theta)\,\Psi_0^S(\rho,\sigma,x)\,. \end{aligned} \tag{23}$$

The corresponding integrated photoabsorption cross section for a single excited electric dipole state, in the long wavelength limit [17, 24], is

$$\Sigma_1 = \int dE_\gamma \sigma_\gamma^{E1} = \frac{4\pi^2\hbar c}{\hbar\omega}\,\overline{\left|\mathcal{M}_{E1}\right|^2}. \tag{24}$$

For the $M1$ transition amplitude we have

$$M_{M1} = < \Psi_f | H_{M1} | \Psi_i > = i\left(\hat{\epsilon}_\gamma \times \vec{k_\gamma}\right) \sum_{j=1}^{3} < \Psi_f | \mu_q^j \sigma_j^z | \Psi_i >\,, \tag{25}$$

where Ψ_i and Ψ_f are given by Eqs. (10) and (12) respectively.

The process $\gamma P \to \Delta^+(1232)$ $(\frac{1}{2}^+ \to \frac{3}{2}^+)$ can take place either via the magnetic dipole ($M1$) or the electric quadrupole ($E2$) transition. In the quark model the latter transition is forbidden [25] because it is proportional to the charge operator which cannot cause transitions between quark spin 1/2 and 3/2 states, and hence the matrix element vanishes by orthogonality of the quark spin wave functions. The $M1$ transition involves the quark magnetic moments – hence the spin operator – and this can lead to transitions $(S = 1/2) \to (S = 3/2)$ [26]. The transition matrix element can be written as

$$M_{M1} = i\left(\hat{\epsilon}_\gamma \times \vec{k_\gamma}\right) < \Psi_f | \mu_q^1 \sigma_1^z + \mu_q^2 \sigma_2^z + \mu_q^3 \sigma_3^z | \Psi_i > \ . \tag{26}$$

Using the the flavor, spin, and configuration space wave functions and averaging over the two photon polarization directions we finally obtain

$$\overline{\left|\mathcal{M}_{M1}\right|^2} = \frac{2\alpha}{9}\frac{E_\gamma^2\hbar c}{(mc^2)^2} I_{M1}^2 \,, \tag{27}$$

where I_{M1} is the overlap integral given by

$$I_{M1} = 8\pi^2 \int_0^\infty \rho^2 d\rho \int_0^\infty \sigma^2 d\sigma \int_{-1}^{1} dx\, \Psi_\Delta^S(\rho,\sigma,x)\Psi_0^S(\rho,\sigma,x)\,. \tag{28}$$

Like in the electric transitions, the photoabsorption cross section for a single excited magnetic dipole state is

$$\Sigma_{M1} = \int dE_\gamma \sigma_\Delta^{M1} = \frac{4\pi^2\hbar c}{\hbar\omega}\,\overline{\left|\mathcal{M}_{M1}\right|^2}\,. \tag{29}$$

4. – Results and Discussions

The quark-quark potential can be written as a sum of the central and the spin-spin parts

$$V_{qq} = V^c + V^s \,, \tag{30}$$

where V^c contains, as usual, the confinement and the coulombic parts while V^s is of the general form

$$V^s = f_{ij}(r)\vec{\sigma}_i\cdot\vec{\sigma}_j\,. \tag{31}$$

The spin-dependent interaction, suggested by the perturbative one-gluon exchange mechanism, is important in describing the splitting of the meson masses and the experimental mass difference $M_\Delta - M_N$. It contains the function $f_{ij}(r)$ which is singular as $\delta(r)$ and therefore the corresponding wave equation has no physically acceptable solutions and one expects a collapse in both the quark-quark and in baryon systems. To avoid this a cut-off or a smearing function is introduced to reduce the singularity and to make the calculations tractable. Since the resulting spin-spin potential is strongly attractive and short-ranged, it can generate significant short–range correlations which have stronger effects on the ground state than on the excited $L > 0$ states. The strong correlations at short distances in the $L = 0$ partial wave are not expected to play a major role in photoexcitation processes as the excited states are shifted to the outer region and thus in the overlap integral short-range contributions are rather unimportant. Therefore any influence of the spin-spin force will come from the modification of the mass differences and of the corresponding wave functions. In this work we performed calculations with the V^c term alone as well as with both terms included. For the calculations with the V^c part only, we employed the Martin [27, 28], the Cornell [29, 30], and the Lichtenberg [31] potentials. For the calculations in which both terms are included we used the Ono-Schöberl potential [32] and two of the recently published potentials by Silvestre-Brac [1], namely, the AP1 and AP2 versions which have the general form

$$\begin{aligned} V_{q\bar{q}}(r_{ij}) = &-\frac{\kappa(1-\exp(-r_{ij}/r_c))}{r_{ij}} + \lambda r_{ij}^p - \Lambda \\ &+ \frac{2\pi}{3m_im_j}\kappa'(1-\exp(-r_{ij}/r_c))\frac{\exp(-r_{ij}^2/r_0^2)}{\pi^{3/2}r_0^3}\vec{\sigma}_i\vec{\sigma}_j \end{aligned} \tag{32}$$

with

$$r_0(m_i, m_j) = A \left(\frac{2m_i m_j}{m_i + m_j} \right)^{-B} .$$

The parameters for AP1 are given by

$$\begin{array}{lll} p = 2/3\,, & r_c = 0\,, & m_u = m_d = 0.277\,\mathrm{GeV}\,, \\ \kappa = 0.4242\,, & \kappa' = 1.8025\,, & \lambda = 0.3898\,\mathrm{GeV}^{5/3}\,, \\ \Lambda = 1.1313\,\mathrm{GeV}\,, & B = 0.3263\,, & A = 1.5296\,\mathrm{GeV}^{B-1}\,, \end{array}$$

while those of AP2 are given by

$$\begin{array}{lll} p = 2/3\,, & r_c = 0.3466\,\mathrm{GeV}^{-1}\,, & m_u = m_d = 0.280\,\mathrm{GeV}\,, \\ \kappa = 0.5743\,, & \kappa' = 1.8993\,, & \lambda = 0.3978\,\mathrm{GeV}^{5/3}\,, \\ \Lambda = 1.1146\,\mathrm{GeV}\,, & B = 0.3478\,, & A = 1.5321\,\mathrm{GeV}^{B-1}\,. \end{array}$$

We remind here the reader that the quark-quark potential V_{qq} is related to the quark-antiquark potential $V_{q\bar{q}}$ by Lipkin's rule [33], i.e., $V_{qq} = V_{q\bar{q}}/2$.

The results obtained for the ground state and the first orbitally excited state of the nucleon with spin-independent potentials and by using the IDEA and HCA methods are given in table I. Both methods are in very good agreement with each other and the transition energy $E_1 - E_0$, shown in table II, is reasonably well reproduced. This transition energy and the corresponding wave functions were used to calculate the integrated photoabsorption cross section Σ_1 due to the E1 transition, from the ground state to the N(1520) resonance. Our results are given in table II together with experimental values and those of other methods. It is seen that for both the IDEA and the HCA the cross sections are in good to very good agreement with the experimental data.

The results obtained by other methods for the transition energy are generally very low as compared to the experimental ones except those obtained via the Isgur-Karl (I.K.) model [18], which, nevertheless, can not reproduce the experimental photoabsorption cross section Σ_1 well. This mainly implies that the corresponding wave functions are not adequate to describe the photoabsorption cross section for the E1 transition. The latter is also true for the wave functions employed by Brizzolara and Giannini in their thorough investigations on the nucleon photoabsorption [17] based on the nonrelativistic quark model of Isgur and Karl and the bag model. They explicitly demonstrated that the results are rather strongly dependent on the hyperfine mixing and on the dimensions of the system and thus on the h.o parameter used. The discrepancy with the experimental results still persists even when other contributions such as charge plus current densities with and without retardation effects are incorporated into the h.o model. Three-body force components of the quark interaction (in the 3q model), and relativistic corrections in the framework of the bag model do not improve matters to a significant degree either. Apart, of course, from the inadequacy of the wave functions these discrepancies probably stem also from some other fundamental absorption mechanism, such as the coupling to mesons and/or $q\bar{q}$ pairs [17].

The results for the nucleon masses obtained by using spin-dependent potentials are given in table III. The corresponding transition energy $E_1 - E_0$ together with the cross sections are presented in table IV. It is seen that both AP1 and AP2 potentials of Silvestre-Brac give excellent results for the photoabsorption cross section Σ_1. This came as a surprise since the transition energies are not as good as one would like them

to be. This can be attributed to the fact that in the confining potential the excited state lies at a higher energy as compared to the experimental value and thus the wave function for $L = 1$ is more spread out in space and therefore the overlap integral somehow compensates for the excess transition energy. Our results obtained with the Ono-Schöberl potential are in fair agreement with the experimental data.

The delta masses are given in table V while the transition energy $E_\Delta - E_N$ is shown in table VI; it is well reproduced but the results for the corresponding integrated photoabsorption cross section of the $\Delta(1232)$ resonance are only in fair agreement with the experimental values. The same is true for the I.K model. This discrepancy can be connected to the strong dependence on the quark masses that enter the cross section via the magnetic moments.

Finally the energy and cross–section results obtained via the IDEA and HCA methods are quite similar. This comes as no major surprise as the quark-quark potentials are soft and thus two-body correlations do not play a significant role. This means that the corresponding wave functions are fairly similar and the small differences between them do not strongly influence the overlap integral and thus the cross sections. In the case where a spin-spin force was used the short–range correlations are not manifested in the overlap integral either due to the shifting of the $L > 0$ wave function to the outer region.

In conclusion the nonrelativistic potential model reproduced the experimental data for the E1 transitions quite successfully. There is a weak potential dependence and in the case of the spin-independent potentials the best results are obtained with the Lichtenberg potential. In the case of the spin-dependent potentials both AP1 and AP2 give excellent results for the E1 photoabsorption cross sections. In the M1 transitions where there is a strong dependence on the quark masses the Ono-Schöberl potential gives better results than the AP1 and AP2 potentials.

* * *

Financial support from the Foundation for Research Development is gratefully acknowledged.

REFERENCES

[1] Silvestre-Brac B., *Few-Body Systems*, **20** (1996) 1.
[2] Lucha W., Schöberl F. F. and Gromes D., *Phys. Rep.*, **200** (1991) 127.
[3] Richard J. -M., *Phys. Rep.*, **212** (1992) 1.
[4] Fabre de la Ripelle M., *Phys. Lett. B*, **205** (1988) 97.
[5] Badalyan A. M., Kitoroage D. I. and Pariyski D. S., *Nucl. Phys.*, **46** (1987) 226.
[6] Badalyan A. M. and Kitoroage D. I., *Nucl. Phys.*, **47** (1987) 807.
[7] Haysak M. I., Lengyel V. I., Chalupka S. and Salak M., *Czech. J. of Physics*, **41** (1991) 799.
[8] Fabre de la Ripelle M., Fiedeldey H. and Sofianos S. A., *Few-Body Systems Suppl.*, **2** (1987) 493.
[9] McTavish J. P., Fiedeldey H., Fabre de la Ripelle M. and van der Merwe P. du T., *Few-Body Systems*, **3** (1988) 99.
[10] Giannini M. M., *Nucl. Phys. A*, **497** (1989) 537c.
[11] Kvitsinsky A. A., Kuperin Yu. A., Merkurev S. P. and Yu Novozhilov V., *Sov. J. Nucl. Phys.*, **39** (1983) 419; Silvestre-Brac B. and Gignoux C., *Phys. Rev. D*, **32** (1985) 743.
[12] Fabre de la Ripelle M., in *Models and Methods in Few-Body Physics*, Vol. 273 of *Lecture Notes in Physics*, edited by Fereira L. S., Fonseca A. C. and Streit L., (Springer-Verlag, New York, 1986), pp. 283-323.
[13] Fabre de la Ripelle M., Fiedeldey H. and Sofianos S. A., *Phys. Rev. C*, **38** (1988) 449; *Few-Body Systems*, **6** (1989) 157; *Few-Body Systems Suppl.*, **2** (1987) 488.
[14] Oehm W., Sofianos S. A., Fiedeldey H. and Fabre de la Ripelle M., *Phys. Rev. C*, **44** (1991) 81.
[15] Oehm W., Sofianos S. A., Fiedeldey H. and Fabre de la Ripelle M., *Phys. Rev. C*, **43** (1991) 25.
[16] Giannini M. M., *Rep. Prog. Phys.*, **54** (1990) 453.
[17] Brizolara R. and Giannini M. M., *Few–body Systems*, **3** (1987) 49; **5** (1988) 1.
[18] Isgur N. and Karl G., *Phys. Rev. D*, **18** (1978) 4187; **19** (1979) 2653.
[19] Close F. E., *An Introduction to Quarks and Partons*, Academic Press, 1979.
[20] Roy R. R. and Nigam B. P., *Nuclear Physics*, New York, John Wiley and Sons 1967.
[21] Siegert A. J. F., *Phys. Rev.*, **52** (1937) 787.
[22] Ellerkmann G., Sandhas W., Sofianos S. A. and Fiedeldey H., *Phys. Rev. C*, **53** (1996) 2638.
[23] Böttger W., Casel A. and Sandhas W., *Phys. Lett. B*, **92** (1980) 11.
[24] Eisenberg J. M., Greiner W., *Excitation Mechanisms of the Nucleus*, Amsterdam, North Holland 1970.
[25] Becchi C. and Morpurgo G., *Phys. Lett.*, **17** (1965) 352.
[26] Dalitz R. H. and Sutherland D., *Phys. Rev.*, **146** (1966) 1180.
[27] Martin A., *Phys. Lett. B*, **100** (1981) 511.
[28] Richard J. -M. and Taxil P., *Phys. Lett. B*, **128** (1983) 453; *Ann. Phys. (N. Y.)*, **150** (1983) 267.
[29] Eichten E., Gottfried K., Kinoshita T., Kogut J., Lane K. D. and Yan T. -M., *Phys. Rev. Lett.*, **34** (1975) 369.
[30] Badalyan A. M., *Phys. Lett. B*, **199** (1987) 267.

[31] Lichtenberg D. B. and Tassie L. J., *Phys. Rev.*, **155** (1967) 1601.
[32] Ono S. and Schöberl F., *Phys. Lett. B*, **118** (1982) 419; Schöberl F., Falkensteiner P. and Ono S., *Phys. Rev. D*, **30** (1984) 603.
[33] Greenberg O. W. and Lipkin H. J., *Nucl. Phys. A*, **370** (1981) 349.
[34] Armstrong T. A. *et al.*, *Nucl. Phys. B*, **41** (1972) 445.

TABLE I. – Nucleon masses in MeV obtained with spin-independent potentials.

Quantum numbers			Expt	IDEA Potentials			HCA Potentials		
L	S	J^P	values	M	C	L	M	C	L
0	1/2	$\frac{1}{2}^+$	940	1085	1087	1085	1090	1090	1088
1	1/2	$\frac{3}{2}^-$	1520	1676	1773	1730	1676	1774	1730

M : Martin potential; C : Cornell potential; L : Lichtenberg potential.

TABLE II. – The integrated photoabsorption cross section Σ_1 for the N(1520) proton resonance obtained with spin-independent potentials. Abbreviations as in table I, see also text.

	Expt	IDEA Potentials			HCA Potentials		
	values	M	C	L	M	C	L
$E_1 - E_0$ (MeV)	580	591	686	645	586	684	642
Σ_1 (MeVmb)	33[a]	30.4	27.8	30.6	30.6	28.6	31.3
			OTHER	METHODS			
		I.K.	h.o.	h.o. (ret)	h.o. $(\rho + j)$	3q	bag
$E_1 - E_0$ (MeV)	580	595[b]	167	167	167	144	349
Σ_1 (MeVmb)	33[a]	26[c]	40[c]	38	43	27	13

[a]Ref. [34]; [b]Ref. [18]; [c]Ref. [17].
The symbols ret, ρ, and j mean retardation, classical charge density and current density respectively.

TABLE III. – Nucleon masses in MeV obtained with spin-dependent potentials.

Quantum numbers				IDEA Potentials			HCA Potentials		
L	S	J^P	Expt values	AP1	AP2	OS	AP1	AP2	OS
0	1/2	$\frac{1}{2}^+$	940	1028	1021	920	1035	1029	940
1	1/2	$\frac{3}{2}^-$	1520	1772	1797	1821	1773	1797	1821

OS : Ono-Schöberl potential; AP1, AP2 : Silvestre-Brac potentials.

TABLE IV. – The integrated photoabsorption cross section Σ_1 for the N(1520) proton resonance obtained with spin-dependent potentials. Abbreviations as in table III, see also text.

		IDEA Potentials			HCA Potentials		
	Expt values	AP1	AP2	OS	AP1	AP2	OS
$E_1 - E_0$ (MeV)	580	744	776	901	738	768	881
Σ_1 (MeVmb)	33[a]	32.6	32.3	27.1	32.9	32.5	27.4

TABLE V. – Delta masses in MeV obtained with spin-dependent potentials. Abbreviations as in table III, see also text.

Quantum numbers				IDEA Potentials			HCA Potentials		
L	S	J^P	Expt values	AP1	AP2	OS	AP1	AP2	OS
0	3/2	$\frac{3}{2}^+$	1232	1300	1307	1224	1300	1308	1225

TABLE VI. – The integrated photoabsorption cross section Σ_{M1} for the $\Delta(1232)$ resonance obtained with spin-dependent potentials. Abbreviations as in tables II and III, see also text.

		IDEA Potentials			HCA Potentials			OTHER METHODS
	Expt values	AP1	AP2	OS	AP1	AP2	OS	I.K.
$E_\Delta - E_N$ (MeV)	292	272	286	304	265	279	285	300[b]
Σ_{M1} (MeVmb)	63[a]	81	83	54	79	82	58	52